

# Kaon effective mass and energy from a novel chiral SU(3)-symmetric Lagrangian

Guangjun Mao<sup>1</sup>, P. Papazoglou<sup>1</sup>, S. Hofmann<sup>1</sup>, S. Schramm<sup>2</sup>, H. Stöcker<sup>1</sup>, and  
W. Greiner<sup>1</sup>

<sup>1</sup>*Institut für Theoretische Physik der J. W. Goethe-Universität*

*Postfach 11 19 32, D-60054 Frankfurt am Main, Germany*

<sup>2</sup>*GSI Darmstadt, Postfach 11 05 52, D-64220 Darmstadt, Germany*

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A new chiral SU(3) Lagrangian is proposed to describe the properties of kaons and antikaons in the nuclear medium, the ground state of dense matter and the kaon-nuclear interactions consistently. The saturation properties of nuclear matter are reproduced as well as the results of the Dirac-Brückner theory. After taking into account the coupling between the omega meson and the kaon, we obtain similar results for the effective kaon and antikaon energies as calculated in the one-boson-exchange model while in our model the parameters of the kaon-nuclear interactions are constrained by the SU(3) chiral symmetry.

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The properties of kaons/antikaons in nuclear and neutron matter have attracted considerable interest since the pioneering work of Kaplan and Nelson [1], who first proposed the occurrence of kaon condensation at several times normal nuclear density. Considerable theoretical effort has been devoted to investigate such medium effects on kaons and antikaons in dense matter [2–9]. Recent data by the Kaos collaboration [10] on  $K^+$  and  $K^-$  production in relativistic heavy-ion collisions, which seems to exhibit a substantial enhancement of the  $K^-$  yield, stimulated further activity [11]. Among the various models proposed, the chiral SU(3) Lagrangian seems to be particularly useful, since the kaon is essentially a pseudo-Goldstone boson. The effective Lagrangian of chiral perturbation theory used in Refs. [1,3] reads

$$\mathcal{L}_{KN} = -\frac{3i}{8f_K^2}\bar{\psi}\gamma_\mu\psi\bar{K}\overset{\leftrightarrow}{\partial}^\mu K + \frac{\Sigma_{KN}}{f_K^2}\bar{\psi}\psi\bar{K}K. \quad (1)$$

Here the iso-spin dependent terms have been dropped since the problem will be discussed in symmetric nuclear matter. The second term of the above equation stems from the explicit breaking of chiral symmetry. Usually, only the linear order of current-quark mass is taken into account, which leads to  $f_K = f_\pi = 93$  MeV [12]; therefore, in the following  $f_K$  in Eq. (1) is replaced by  $f_\pi$ . The amplitude of the  $KN$  sigma term has not yet been determined unambiguously. Two different choices,  $\Sigma_{KN} \approx 2m_\pi$  and  $\Sigma_{KN} = 450 \pm 30$  MeV, have been proposed, in accordance with the Bonn model [2] and with lattice gauge calculations [13], respectively. The kaon and antikaon effective-masses and -energies in static nuclear matter can be easily derived from Eq. (1)

$$m_K^{*2} = m_K^2 - \frac{\Sigma_{KN}}{f_\pi^2}\rho_S, \quad (2)$$

$$\omega_{K,\bar{K}} = \sqrt{m_K^{*2} + \left(\frac{3}{8f_\pi^2}\rho_B\right)^2} \pm \frac{3}{8f_\pi^2}\rho_B. \quad (3)$$

The minus sign in Eq. (3) corresponds to the antikaon energy.  $\rho_S$  and  $\rho_B$  are the scalar and the vector (net) baryon density, respectively. Since chiral perturbation theory has no direct relation to the ground state properties of the dense matter, one usually uses [5] the  $\rho_S$  and  $\rho_B$  vacuous as calculated by the relativistic mean-field theory of Walecka model

[14], i.e., a non-chiral Lagrangian, or simply set  $\rho_S = \rho_B$  [1], in order to evaluate Eqs. (2) and (3). Therefore, the self-consistency of the theory is lost.

A different approach, based on the chiral SU(3) Lagrangian, is the coupled channel model [8,9]. By including the  $\Lambda(1405)$  as a  $K^-p$  quasi-bound state, the model gives a strong, non-linear density dependence of the  $K^-$  potential, which changes sign from positive to negative values at low density (around  $0.1\rho_0$ , where  $\rho_0$  is the ground-state nuclear matter density). This trick allows for an attractive potential of antikaons without violating the low density theorems. Once the density exceeds  $0.2\rho_0$ , the repulsive effect of the  $\Lambda(1405)$  is neglected - it is predicted that the  $\Lambda(1405)$  melts in the dense medium, in analogy to a Mott phase transition (for an alternative analysis see, however, Ref. [15]). However, this model does not take into account the saturation properties of the system. Both pion, nucleon and hyperon masses are kept constant in these calculations, at different densities. Up to now, a consistent calculation based on a chiral Lagrangian, which can simultaneously describe both the kaon-nuclear interactions and the ground state of the dense matter, has not been performed yet.

This is the aim of the present work. It addresses these problems in a novel chiral SU(3)-symmetric Lagrangian [16]. In addition to the ground-state saturation properties of nuclear matter, the whole density dependence of the mean fields as predicted by the Dirac-Brückner theory [17] are considered as a further constraint to the model. This will turn out to be rather important for the investigation of the kaon and antikaon properties at higher densities. The chiral SU(3) Lagrangian reads

$$\mathcal{L} = \mathcal{L}_{kin} + \sum_{W=X,V,u,\Gamma,A} \mathcal{L}_{BW} + \mathcal{L}_{vec} + \mathcal{L}_0 + \mathcal{L}_{SB}. \quad (4)$$

Here  $\mathcal{L}_{kin}$  is the kinetic energy term,  $\mathcal{L}_{BW}$  gives the various interaction terms of the different baryons with 4 lowest (spin-0 and spin-1) meson-multiples and with the photons.  $\mathcal{L}_{vec}$  generates the vector meson-masses through interactions with the spin-0 mesons, and  $\mathcal{L}_0$  gives the meson-meson interaction potentials which includes the spontaneous breaking of chiral symmetry and trace anomaly. Finally,  $\mathcal{L}_{SB}$  introduces an explicit symmetry breaking of  $U(1)_A$ ,  $SU(3)_V$ , and the chiral symmetry. The main feature of the model is

that the baryon masses are generated by the scalar condensates while their splitting is realized through SU(3) symmetry breaking for these condensates. The model is described in [16]. Considering SU(3) generators up to quadratic order, the Lagrangian for nuclear matter reads

$$\mathcal{L}_{kin} = i\bar{\psi}\gamma_\mu\partial^\mu\psi + \frac{1}{2}\partial_\mu\sigma\partial^\mu\sigma + \frac{1}{2}\partial_\mu\zeta\partial^\mu\zeta + \frac{1}{2}\partial_\mu\chi\partial^\mu\chi - \frac{1}{4}\omega_{\mu\nu}\omega^{\mu\nu}, \quad (5)$$

$$\mathcal{L}_{NX} + \mathcal{L}_{NV} = -\bar{\psi} [g_{N\omega}\gamma_\mu\omega^\mu + m_N] \psi, \quad (6)$$

$$\mathcal{L}_{vec} = \frac{1}{2}m_\omega^2\frac{\chi^2}{\chi_0^2}\omega_\mu\omega^\mu + g_4^4(\omega_\mu\omega^\mu)^2, \quad (7)$$

$$\begin{aligned} \mathcal{L}_0 = & -\frac{1}{2}k_0\chi^2(\sigma^2 + \zeta^2) + k_1(\sigma^2 + \zeta^2)^2 + k_2\left(\frac{\sigma^4}{2} + \zeta^4\right) + k_3\chi\sigma^2\zeta \\ & -k_4\chi^4 + \frac{1}{4}\chi^4\ln\frac{\chi^4}{\chi_0^4} + \frac{\delta}{3}\chi^4\ln\frac{\sigma^2\zeta}{\sigma_0^2\zeta_0}, \end{aligned} \quad (8)$$

$$\mathcal{L}_{SB} = -\left(\frac{\chi}{\chi_0}\right)^2 (f_1\sigma + f_2\zeta), \quad (9)$$

and the kaon interaction is described by

$$\mathcal{L}_{KN} = -\frac{3i}{8f_K^2}\bar{\psi}\gamma_\mu\psi\bar{K}\overset{\leftrightarrow}{\partial}^\mu K + \frac{m_K^2}{2f_K}(\sigma + \sqrt{2}\zeta)\bar{K}K - ig_{\omega K}\bar{K}\overset{\leftrightarrow}{\partial}^\mu K\omega_\mu. \quad (10)$$

Here  $\sigma$ ,  $\omega$  are the scalar and vector field and  $\zeta$ ,  $\chi$  are the strange scalar field and the gluon field, respectively;  $f_1 = m_\pi^2 f_\pi$ ,  $f_2 = \sqrt{2}m_K^2 f_K - \frac{1}{\sqrt{2}}m_\pi^2 f_\pi$ ,  $m_N = m_0 - \frac{1}{3}g_8^S(4\alpha - 1)(\sqrt{2}\zeta - \sigma)$ ,  $m_0 = g_1^S(\sqrt{\frac{2}{3}}\sigma + \sqrt{\frac{1}{3}}\zeta)$ . The omega-kaon coupling is introduced through considering the vector field as a gauge field. The vacuum condensates of the scalar fields  $\sigma_0$  and  $\zeta_0$  generate the masses of the various hadrons. They are constrained by the pion and kaon decay constants:  $f_\pi = -\sigma_0$ ,  $f_K = -(\sigma_0 + \sqrt{2}\zeta_0)/2$ . If the equality  $\sigma_0 = \sqrt{2}\zeta_0$  is satisfied, the model regains the SU(3) symmetry. The parameters of the model are  $\sigma_0$ ,  $\zeta_0$ ,  $\chi_0$ ,  $g_{N\omega}$ ,  $\sum_{i=0}^4 k_i$ ,  $\delta$ ,  $g_4$  and  $g_1^S$ ,  $g_8^S$ ,  $\alpha$ . Ten of these are determined by the SU(3) vacuum and the 18+8 hadron masses, three parameters, i.e., the vector coupling constant  $g_{N\omega}$  and  $g_4$  plus the "gluon condensate"  $\chi_0$  are used to fit the saturation properties of nuclear matter. This yields two sets of parameters denoted as C1 and C2. These two parameters sets differ in the strange condensate  $\zeta$ . C1 allows for an explicit  $\zeta$ -dependence of  $m_N$ , while C2 excludes such a dependence. The corresponding saturation properties are (1)

C1:  $m^*/m_N = 0.612$ ,  $E/A(\rho_0) = -15.99$  MeV,  $K = 276.3$  MeV; (2) C2:  $m^*/m_N = 0.641$ ,  $E/A(\rho_0) = -15.93$  MeV,  $K = 266.2$  MeV. Both C1 and C2 have a saturation density  $\rho_0 = 0.15 \text{ fm}^{-3}$  and  $f_\pi = 93.3$  MeV,  $f_K = 122$  MeV,  $m_\pi = 139$  MeV,  $m_K = 498$  MeV. The parameters of the kaon-nuclear interactions, Eq. (10), are constrained by the chiral Lagrangian itself.  $g_{\omega K} = g_{\rho\pi\pi}f_\pi^2/2f_K^2$ , and  $g_{\rho\pi\pi} = 6.05$  from the  $\rho^0 \rightarrow \pi^+\pi^-$  decay [18]. The results do not depend on the  $KN$  sigma term.  $\Sigma_{KN}$  is computed in the present model [19].

A field shift to new variables,  $\phi$  and  $\xi$ , is performed ( $\sigma = \sigma_0 - \phi$ ,  $\zeta = \zeta_0 - \xi$ ) and  $\chi = \chi_0$  is set (the variation of the gluon condensate in the nuclear medium is negligible [16]). Then the following field equations of the scalar and vector mesons in static nuclear matter are obtained after some straightforward algebra

$$\begin{aligned} & (k_0\chi_0^2 - 12k_1\sigma_0^2 - 4k_1\zeta_0^2 - 6k_2\sigma_0^2 - 2k_3\chi_0\zeta_0)\phi + (12k_1\sigma_0 + 6k_2\sigma_0)\phi^2 \\ & - (4k_1 + 2k_2)\phi^3 + \frac{2\delta}{3}\chi_0^4\frac{1}{\sigma_0 - \phi} - \frac{2\delta}{3\sigma_0}\chi_0^4 - 4k_1\phi\xi^2 + (8k_1\zeta_0 + 2k_3\chi_0)\phi\xi \\ & + 4k_1\sigma_0\xi^2 - (8k_1\sigma_0\zeta_0 + 2k_3\chi_0\sigma_0)\xi = -g_{N\sigma}\rho_S, \end{aligned} \quad (11)$$

$$\begin{aligned} & (k_0\chi_0^2 - 12k_1\zeta_0^2 - 4k_1\sigma_0^2 - 12k_2\zeta_0^2)\xi + 12(k_1 + k_2)\zeta_0\xi^2 - 4(k_1 + k_2)\xi^3 \\ & + \frac{\delta}{3}\chi_0^4\frac{1}{\zeta_0 - \xi} - \frac{\delta}{3\zeta_0}\chi_0^4 - 4k_1\phi^2\xi + (4k_1\zeta_0 + k_3\chi_0)\phi^2 + 8k_1\sigma_0\phi\xi \\ & - (8k_1\sigma_0\zeta_0 + 2k_3\chi_0\sigma_0)\phi = -g_{N\zeta}\rho_S, \end{aligned} \quad (12)$$

$$m_\omega^2\omega + 4g_4^4\omega^3 = g_{N\omega}\rho_B. \quad (13)$$

The effective-mass and -energy of the kaon  $K$  and the antikaon  $\bar{K}$  are given by

$$m_K^{*2} = m_K^2 + \frac{m_K^2}{2f_K}\phi + \frac{m_K^2}{\sqrt{2}f_K}\xi, \quad (14)$$

$$\omega_{K,\bar{K}} = \sqrt{m_K^{*2} + \left(\frac{3}{8f_K^2}\rho_B + g_{\omega K}\omega_0\right)^2} \pm \left(\frac{3}{8f_K^2}\rho_B + g_{\omega K}\omega_0\right). \quad (15)$$

Here the coupling strengths  $g_{N\sigma}$  and  $g_{N\zeta}$  are given by

$$g_{N\sigma} = - \left[ \sqrt{\frac{2}{3}}g_1^S + \frac{1}{3}g_8^S(4\alpha - 1) \right], \quad (16)$$

$$g_{N\zeta} = - \left[ \sqrt{\frac{1}{3}}g_1^S - \frac{\sqrt{2}}{3}g_8^S(4\alpha - 1) \right]. \quad (17)$$

In principle, one can implement additional terms on the Lagrangian with more than one time derivate acting on the kaon field (i.e., the so-called off-shell terms), for example,  $\sim Tr(u_\mu u^\mu \bar{B}B)$ . These terms are not included in the present work.

Fig. 1 displays the binding energy of the system (a) and the scalar and vector potentials of the nucleons (b) as a function of the Fermi momentum. The results of the chiral effective Lagrangian with the C1 and C2 parameters as well as the relativistic mean field theory with the linear [14] and non-linear (TM1) [20] parameterization are compared with the prediction of the Dirac-Brückner G-matrix theory [17]. It can be seen that the results of the linear Walecka model deviate from the Dirac-Brückner theory evidently. However, C1, C2 and TM1 can reproduce the results of the relativistic G-matrix theory nearly perfectly up to the normal density. This explains why the description of finite nuclei is rather convincing in these effective models [20,16]. At higher densities, the results of C1 and C2 are closer to the prediction of the Dirac-Brückner theory than TM1. Both C1 and C2 follow the results of the G-matrix calculations closely up to four times normal density although the parameters have not been fitted to the results of Ref. [17]. Therefore, the novel chiral SU(3) Lagrangian [16] yields reasonable results for dense matter  $\rho \leq 4\rho_0$ . These results provide a sound self-consistent basis for the investigation of kaon and antikaon properties in the nuclear medium.

Fig. 2(a) depicts the kaon and antikaon effective-masses as calculated from this SU(3) chiral model with the parameter set C2 (the results of parameter set C1 are nearly indistinguishable from that of C2). The results of the Kaplan-Nelson model are also plotted for two different  $KN$  sigma terms,  $\Sigma_{KN} = 2m_\pi$  and  $\Sigma_{KN} = 450$  MeV with the mean fields provided by Eqs. (11)-(13) (parameter set C2). The density dependence of the scalar fields as computed with the parameter set C2 is given in Fig. 2(b). From now on, let us denote the Lagrangian of Eq. (1) model I and that of Eq. (4) model II, respectively. For model I, the effective mass of the kaon decreases nearly linearly with increasing density, with a slope determined by  $\Sigma_{KN}$ . For model II  $m_K^*$  first decreases with increasing density, but then it approaches saturation, consistent with the finding of the one-boson-exchange (OBE) model [5]. One can easily understand the different behavior of two models from

Eqs. (2) and (14): In Eq. (2) the  $m_K^*$  depends linearly on  $\rho_S$  (which is, in turn, approximately equal to  $\rho_B$ ). In Eq. (14), the  $m_K^*$  is related to the scalar fields. To obtain the Dirac-Brückner results displayed in Fig. 1, a highly non-linear relation between the  $\phi$ ,  $\xi$  and the  $\rho_S$  is asked for, as can be seen from Eqs. (11) and (12). Fig. 2(b) shows the saturation of the scalar fields at high densities. Consequently, the  $m_K^*$  saturates after  $\rho_B \approx 2\rho_0$ , it never tends to zero. Neither kaon nor antikaon condensation occur in the chiral SU(3) model. In fact, antikaon condensation could occur only if the condensates of the scalar fields,  $\sigma$  and  $\zeta$ , would vanish in the medium. This would constitute a chiral phase transition. Note that antikaon condensation may occur in model I depending on the values of  $\Sigma_{KN}$  as well as on the mean field [5] used in the calculations.

Fig. 3 compares the chiral SU(3) model calculations (with the parameter set C2) for the kaon and antikaon effective-energies, with (solid lines) and without (dot-dashed lines) the omega-kaon coupling, with the results of other models [1,5]. This is done up to  $3\rho_0$ , where all these models should be least unreliable (than in the regime of extremely high densities). In addition, an "empirical" kaon dispersion relation [11] is also presented - it resulted from "fitting" the Kaos data [10] of  $K^+$  and  $K^-$  production in heavy-ion collisions by means of a relativistic transport model [22] by adjusting the real part of the kaon and antikaon optical potential, but ignoring the imaginary part (i.e., the in-medium scattering cross sections and hyperon resonances). It can be seen that without the  $\omega - K$  coupling our model predicts a rather weak potential for the antikaon compared to the predictions of other models. After introducing the  $\omega - K$  coupling, the calculated effective energies for kaon and antikaon are quite similar as obtained in the one-boson-exchange model [5]. The effective kaon and antikaon energies are directly related to their optical potentials. The chiral SU(3) model gives  $U_{opt}^{\bar{K}} = -98.7$  MeV and  $U_{opt}^K = 21.5$  MeV at normal density. The predicted  $K^-$  optical potential is in accordance with the results of model I and OBE [5] (between  $-70$  and  $-100$  MeV). It is close to the empirical value obtained by the standard fit of  $K^-$ -atomic data, which gives  $U_{opt}^{\bar{K}} = -85$  MeV, but much weaker as compared to the non-linear fit, which gives  $U_{opt}^{\bar{K}} = -200 \pm 20$  MeV [21]. At present, there exist no firmly established empirical value for the  $U_{opt}^K$ . An estimate based on impulse approximation

[4,5] gives  $U_{opt}^K \approx 29$  MeV. Our results are in agreement with it.

In summary, we have employed a recent developed chiral SU(3) Lagrangian to investigate the properties of kaons and antikaons in the nuclear medium. The kaon-nuclear interactions and the ground state of the dense matter are described consistently for the first time within a chiral approach. The parameters of the kaon-nuclear interactions are constrained by the SU(3) chiral symmetry. The saturation properties of nuclear matter as well as the results of the Dirac-Brückner theory are well reproduced. Due to the highly non-linear kaon interaction with respect to the density, the kaon/antikaon effective mass is changed only moderately in the nuclear medium. After introducing the coupling between the omega meson and the kaon, we obtain similar results for the kaon and antikaon effective energies as calculated in the one-boson-exchange model, i.e., the kaon feels a weak repulsive potential while the antikaon suffers a strong attractive potential.

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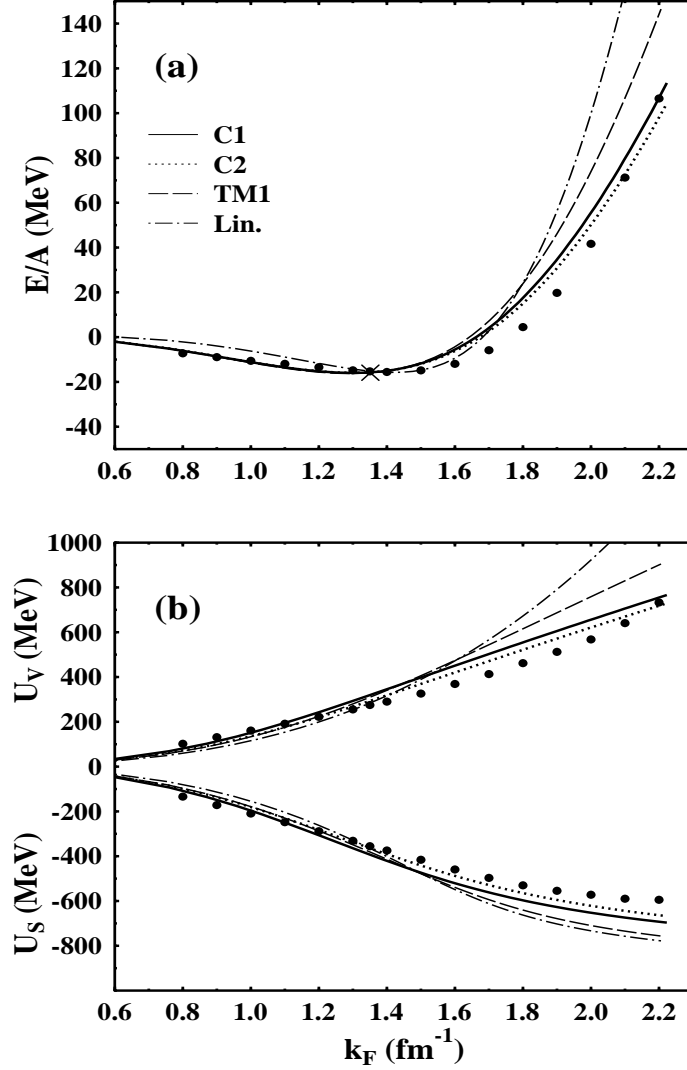


FIG. 1. (a) The binding energy and (b) the scalar and vector potentials of the nucleons calculated with the chiral Lagrangian (C1, C2) and the linear [14] and non-linear (parameter set TM1 [20]) relativistic mean field theory. The dots are the results of the Dirac-Brückner theory of Ref. [17]. The cross in (a) denotes the empirical value for nuclear matter saturation ( $E/A = -16 \pm 1 \text{ MeV}$ ,  $k_F = 1.35 \pm 0.05 \text{ fm}^{-1}$ ). The scalar and vector potential at saturation point ( $k_F = 1.305 \text{ fm}^{-1}$ ) calculated with the chiral Lagrangian are  $-364.3 \text{ MeV}$ ,  $293.3 \text{ MeV}$  for C1 and  $-337.1 \text{ MeV}$ ,  $268.4 \text{ MeV}$  for C2, respectively.

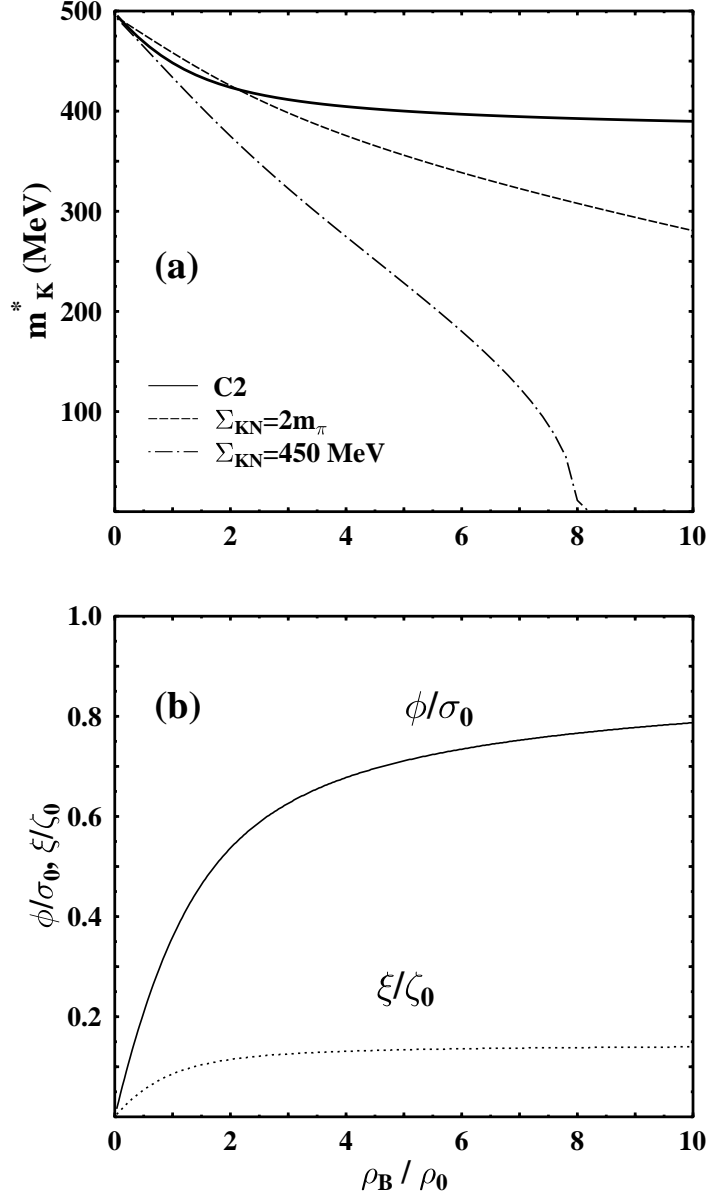


FIG. 2. (a) The effective masses of kaons and antikaons in nuclear matter versus the baryon density at  $T=0$ . The solid curve represents the results of this work with the parameter set C2. The dashed and dot-dashed curves are calculated from Eqs. (2) and (3) with the  $\rho_S$ - and  $\rho_B$ -values as provided by Eqs. (11) - (13) (parameter set C2). (b) The density dependence of the scalar fields as shown for the parameter set C2.

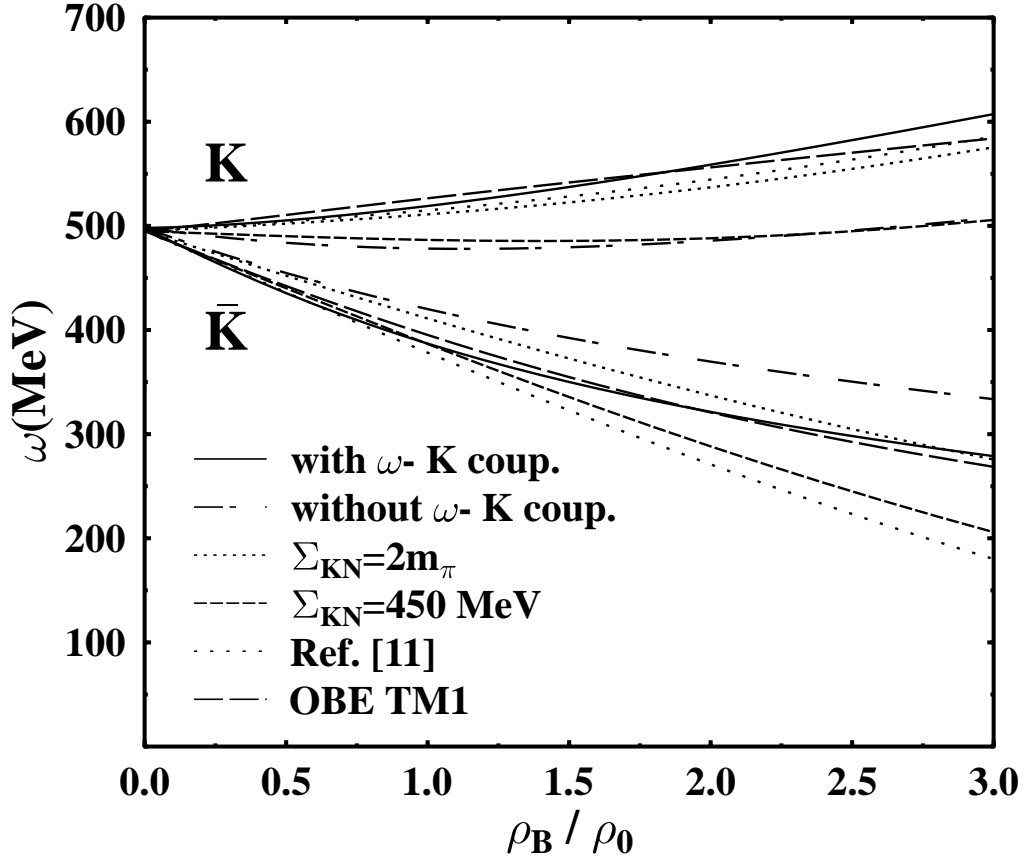


FIG. 3. The energies of kaons and antikaons as a function of the density. The solid and dot-dashed lines represent the results of this work with and without the  $\omega - K$  coupling, respectively. The short-dotted and short-dashed lines are calculated with the chiral perturbation theory with the different  $KN$  sigma terms. The long-dotted lines are the results of Ref. [11]. The long-dashed lines depict the results of the one-boson-exchange model with the TM1 parameter set.